

Unveiling secret interactions among sterile neutrinos with big-bang nucleosynthesis

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(Dated: November 18, 2014)

Short-baseline neutrino anomalies suggest the existence of low-mass ($m \sim \mathcal{O}(1)$ eV) sterile neutrinos ν_s . These would be efficiently produced in the early universe by oscillations with active neutrino species, leading to a thermal population of the sterile states seemingly incompatible with cosmological observations. In order to relieve this tension it has been recently speculated that new “secret” interactions among sterile neutrinos, mediated by a massive gauge boson X (with $M_X \ll M_W$), can inhibit or suppress the sterile neutrino thermalization, due to the production of a large matter potential term. We note however, that they also generate strong collisional terms in the sterile neutrino sector that induce an efficient sterile neutrino production after a resonance in matter is encountered, increasing their contribution to the number of relativistic particle species N_{eff} . Moreover, for values of the parameters of the ν_s - ν_s interaction for which the resonance takes place at temperature $T \lesssim$ few MeV, significant distortions are produced in the electron (anti)neutrino spectra, altering the abundance of light element in Big Bang Nucleosynthesis (BBN). Using the present determination of ^4He and deuterium primordial abundances we determine the BBN constraints on the model parameters. We find that $^2\text{H}/\text{H}$ density ratio exclude much of the parameter space if one assume a baryon density at the best fit value of Planck experiment, $\Omega_B h^2 = 0.02207$, while bounds become weaker for a higher $\Omega_B h^2 = 0.02261$, the 95 % C.L. upper bound of Planck. Due to the large error on its experimental determination, the helium mass fraction Y_p gives no significant bounds.

PACS numbers: 14.60.St, 14.60.Pq, 98.80.-k 26.35.+c

I. INTRODUCTION

In recent years hints for the existence of (sub)-eV sterile neutrinos ν_s , mixing with the three active species, emerged from different short-baseline neutrino oscillation experiments. Notably, the $\bar{\nu}_\mu \rightarrow \bar{\nu}_e$ oscillations in LSND [1] and MiniBoone [2] experiments (recently constrained by the ICARUS experiment [3]), the $\bar{\nu}_e$ and ν_e disappearance revealed by the Reactor Anomaly [4], and the Gallium Anomaly [5] (see [6] for a recent review) have been described in scenarios with one (dubbed “3+1”) or two (“3+2”) sterile neutrinos (see [7, 8] for the latest analysis).

Low-mass sterile neutrinos would be produced also in the early universe via oscillations of active neutrinos [9]. The preferred range of masses and mixing angles from the laboratory anomalies would lead to a complete thermalization of ν_s ’s [10], in disagreement with recent cosmological analysis [11–18]. Indeed, for sterile neutrino masses of 1 eV, or larger, tight constraints are obtained from structure formation, affected at small-scales by the presence of a fully thermalized massive sterile neutrino species [19]. Furthermore, Big Bang Nucleosynthesis (BBN) marginally allows for a completely thermalized sterile neutrino [20, 21].

In order to reconcile the eV sterile neutrino interpre-

tation of the short-baseline anomalies with the cosmological observations, the most straightforward possibility would be to suppress the sterile neutrino thermalization. One of the proposed mechanism [22] (see also [23]) is to consider a primordial asymmetry L_ν between active neutrinos and antineutrinos. This would add an additional “matter potential” term in the active-sterile neutrino equations of motion. If sufficiently large, this term inhibits the active-sterile flavor conversions via the in-medium suppression of the mixing angle. In recent papers [10, 24, 25] it has been shown that in order to achieve a sufficient suppression of the sterile neutrino abundance, an asymmetry $L_\nu \gtrsim 10^{-2}$ is required. However, such a large value does not seem very natural and moreover, it leads to non-trivial effects on BBN due to the distortions induced on active neutrino spectra, as shown in [25].

More recently, in [26, 27] an alternative method to suppress the sterile neutrino production has been proposed, based on the introduction of new “secret” self-interactions among sterile neutrinos, mediated by a massive gauge boson X , with $M_X \ll M_W$. As in the case of neutrino asymmetries, the self-interactions would generate a matter potential in the flavor evolution equations which suppresses the sterile neutrino abundance. If the new interaction mediator X also couples to dark matter, this might also possibly relieve some of the small-

scale structure problems associated with cold dark matter [27, 28] (nonstandard interactions were also introduced to alleviate problems related to cold dark matter in [29–34] and in the references therein).

Since these new interactions involve only the sterile neutrino sector, they would evade existing limits on secret interactions among active neutrinos [35–37], and therefore seem apparently unconstrained. However, as we will discuss in this paper, these interactions can be unveiled by exploiting cosmological observations. Indeed, when the “secret” matter potential becomes of the order of the active-sterile vacuum oscillation frequency, a resonance is encountered maximizing the in-medium mixing angle [38]. This would lead to Mikheyev-Smirnov-Wolfenstein (MSW)-like resonant flavor conversions among active and sterile neutrinos [38]. Moreover, the presence of strong collisional effects in the sterile neutrino sector, again due to the secret interactions, on one side would damp the MSW conversions, on the other would *enhance* the sterile neutrino production via non-resonant processes [39]. In particular, the latter would bring the active-sterile neutrino ensemble towards the flavor equilibrium [40, 41].

In [26] the authors have performed different multi-momentum simulations of the active-sterile flavor evolution in a simplified scheme involving only one active and one sterile species. They found that for values of the ν_s - ν_s coupling $g_X \gtrsim 10^{-2}$ and masses $M_X \gtrsim 10$ MeV, the resonance can produce an increase in the effective neutrino species N_{eff} , introduced as usual as a way to parameterize the energy density ρ_{rel} of relativistic species at some relevant epoch, in terms of photon contribution ρ_γ

$$\rho_{\text{rel}} = \rho_\gamma \left(1 + \frac{7}{8} \left(\frac{4}{11} \right)^{4/3} N_{\text{eff}} \right). \quad (1)$$

Moreover, resonances occurring at $T \lesssim \text{few MeV}$ happen so late that significant distortions are produced in the electron (anti)neutrino spectra. Both these effects have a potential relevance for the abundance of light elements BBN.

Motivated by these warnings, we investigate in details these effects to obtain constraints on the secret ν_s - ν_s interactions parameter space. As in [26], we work in a situation in which active-sterile flavor conversions would occur at a temperature $T \ll M_X$. This implies that one can reduce the interaction to an effective four-fermion point-like structure, with strength ¹

$$G_X = \frac{\sqrt{2}}{8} \frac{g_X^2}{M_X^2}. \quad (2)$$

¹ The numerical factor $\sqrt{2}/8$ has been included in order to exactly mimic the relation between the Fermi constant G_F , the $SU(2)$ coupling constant g and the W -mass in the Standard Model.

Differently from [26] we will work in a (2+1) scenario that allows to describe more realistically the flavor dynamics. However, in this case computing reliably the spectral distortions and N_{eff} as functions of the secret interaction parameters is a very challenging task, involving time consuming numerical calculations for the flavor evolution. Therefore, we will apply an averaged-momentum approximation, where all neutrinos share the mean thermal momentum. In this limit, the information on the active neutrino distribution distortion is contained in a single (time evolving) parameter, which weights the usual Fermi-Dirac distribution. In other terms, neutrinos will be characterized by a gray-body distribution.

This article is structured as follows: in Section II we present the formalism to study the flavor conversions of active-sterile neutrinos in the presence of secret ν_s - ν_s self-interactions and we show some examples of the flavor evolution. We present also our results for the value of N_{eff} as function of G_X and g_X , the parameters characterizing the strength of the new interaction. In Section III we discuss the impact on BBN, in particular on primordial abundances of ^4He and ^2H and discuss the constraints that present experimental data on these two nuclei yields put on the secret interaction scenario. Finally, in Section IV we conclude.

II. SETUP OF THE FLAVOR EVOLUTION

A. Equations of motion

In this Section we summarize the equations of motion for the (2+1) active-sterile neutrino system in the early universe, using the same notation of [24], to which we address the reader for details. In order to take into account the interplay between oscillations and collisions of neutrinos, it is necessary to describe the neutrino (antineutrino) system in terms of 3×3 density matrices ϱ ($\bar{\varrho}$)²

$$\varrho_{\mathbf{p}} = \begin{pmatrix} \varrho_{ee} & \varrho_{e\mu} & \varrho_{es} \\ \varrho_{\mu e} & \varrho_{\mu\mu} & \varrho_{\mu s} \\ \varrho_{se} & \varrho_{s\mu} & \varrho_{ss} \end{pmatrix}. \quad (3)$$

Since our aim is to perform an extensive scan of the parameter space of the ν_s - ν_s secret interactions, in order to carry out a more treatable numerical analysis, we will consider the averaged-momentum approximation, based on the ansatz $\varrho_{\mathbf{p}}(T) \rightarrow f_{FD}(p) \rho(T)$ (see [24]), where $\rho(T)$ is the density matrix for the mean thermal momentum $\langle p \rangle = 3.15 T$, and $f_{FD}(p)$ is the Fermi-Dirac neutrino equilibrium distribution, and similarly for antineutrinos.

² Here ν_μ refers generically to a non-electron active flavor state.

The evolution equation for the momentum-averaged density matrix ρ is the following [24, 42, 43]

$$i \frac{d\rho}{dt} = [\Omega, \rho] + C[\rho], \quad (4)$$

and a similar expression holds for the antineutrino matrix $\bar{\rho}$. The evolution in terms of the comoving observer proper time t can be easily recast in function of the photon temperature T (see [24] for a detailed treatment). The first term on the right-hand side of Eq. (4) describes the flavor oscillations Hamiltonian, given by

$$\begin{aligned} \Omega = & \frac{M^2}{2} \left\langle \frac{1}{p} \right\rangle + \sqrt{2} G_F \left[-\frac{8\langle p \rangle}{3} \left(\frac{E_\ell}{M_W^2} + \frac{E_\nu}{M_Z^2} \right) + N_\nu \right] \\ & + \sqrt{2} G_X \left[-\frac{8\langle p \rangle E_s}{3M_X^2} + N_s \right], \end{aligned} \quad (5)$$

where $M^2 = \mathcal{U}^\dagger \mathcal{M}^2 \mathcal{U}$ is the neutrino mass matrix. Here $\mathcal{U} = \mathcal{U}(\theta_{e\mu}, \theta_{es}, \theta_{\mu s})$ is the 3×3 active-sterile mixing matrix, parametrized as in [24]. We assume $\theta_{e\mu}$ equal to the active 1 – 3 mixing angle θ_{13} [44], while we fix the active-sterile mixing angles to the best-fit values of the different anomalies [7], namely

$$\sin^2 \theta_{e\mu} = 0.023, \quad (6)$$

$$\sin^2 \theta_{es} = 0.033, \quad (7)$$

$$\sin^2 \theta_{\mu s} = 0.012. \quad (8)$$

The mass-squared matrix

$$\mathcal{M}^2 = \text{diag}(-\Delta m_{\text{atm}}^2/2, +\Delta m_{\text{atm}}^2/2, \Delta m_{\text{st}}^2) \quad (9)$$

is parametrized in terms of the atmospheric mass-squared difference $\Delta m_{\text{atm}}^2 = 2.43 \times 10^{-3} \text{ eV}^2$ [44] and of the active-sterile mass splitting $\Delta m_{\text{st}}^2 = 1.6 \text{ eV}^2$, fixed from the short-baseline fit in 3+1 model [7]. In the following we assume the normal mass hierarchy $\Delta m_{\text{atm}}^2 > 0$.

The terms proportional to the Fermi constant G_F in Eq. (5) encode the standard matter effects in the neutrino oscillations. In particular, the term E_ℓ is related to the energy density of e^\pm pairs, E_ν to the energy density of ν and $\bar{\nu}$, and N_ν is the $\nu - \nu$ interaction term proportional to the neutrino asymmetry. The term proportional to G_X represents the new matter secret potential where E_s is the energy density of ν_s and $\bar{\nu}_s$ ³. This refractive term can induce a MSW-like resonance between the active and sterile states when it becomes of the same order of the vacuum frequency associated with the active-sterile mass-squared splitting [26, 27]. Finally, N_s is the self-interaction term proportional to the sterile neutrino asymmetry. In the following, we will consider the most

conservative scenario, with zero neutrino asymmetries in both the active and sterile sectors, so that $\bar{\rho} = \rho$.

The last term in the right-hand side of Eq. (4) is the collisional term. Following [24], for the Standard Model interactions we write it as

$$\begin{aligned} C_{\text{SM}}[\rho] = & -\frac{i}{2} G_F^2 (\{S^2, \rho - I\} - 2S(\rho - I)S \\ & + \{A^2, (\rho - I)\} + 2A(\bar{\rho} - I)A). \end{aligned} \quad (10)$$

In flavor space, $S = \text{diag}(g_s^e, g_s^\mu, 0)$ and $A = \text{diag}(g_a^e, g_a^\mu, 0)$ with the numerical coefficients for the scattering and annihilation processes of the different flavors (see [24] for the numerical values). Concerning the collisional term associated with the secret interactions, one should change $G_F \rightarrow G_X$ in Eq. (10), and introduce two new matrices of coefficients A_X and S_X for annihilations and scatterings mediated by the gauge boson X , respectively. Since we consider here masses of the new boson larger than MeV, in the relevant temperature range $M_X \gg T$ we can neglect annihilation processes and assume $A_X = 0$. For the scattering matrix we have $S_X = \text{diag}(0, 0, 1)$. Therefore, we get the following expression

$$C_X[\rho] = -\frac{i}{2} G_X^2 (\{S_X^2, \rho - I\} - 2S_X(\rho - I)S_X). \quad (11)$$

The strong collisional effects produce a damping of the resonant transitions and would bring the system towards the flavor equilibrium among the different neutrino species with a rate $\Gamma_t \sim \sin^2 2\theta_m G_X^2$ [39–41], where θ_m is the in-medium mixing [45]. At resonance $\theta_m \simeq \pi/4$, so that both the resonant and non-resonant ν_s production are maximized.

To conclude, we mention that in our average-momentum treatment we cannot account for a redistribution in the energy spectra of ν_s associated with the elastic scatterings, considered in the multi-momentum treatment of [26]. However, this effect would not be expected to produce a major impact on our results, described in Section III. Though more refined bounds on the allowed values of g_X and G_X could be only obtained with a full multi-momentum analysis, the main conclusion, basically that there is a strong tension between the secret interaction scenario and primordial deuterium yield, would not change drastically. Finally, we also stress that a precision computation seems also premature and illusory, given the dependence from unknown or poorly constrained parameters in the active sterile mixing from short-baseline neutrino anomalies such as, for example the $\nu_s - \nu_\tau$ mixing.

B. Sterile neutrino production

In Fig. 1 we show the behavior of the different neutrino refractive and collisional rates normalized to the Hubble rate, versus photon temperature T (see [24] for details).

³ See [27] for an explicit calculation of the neutrino potential associated with the secret interactions.

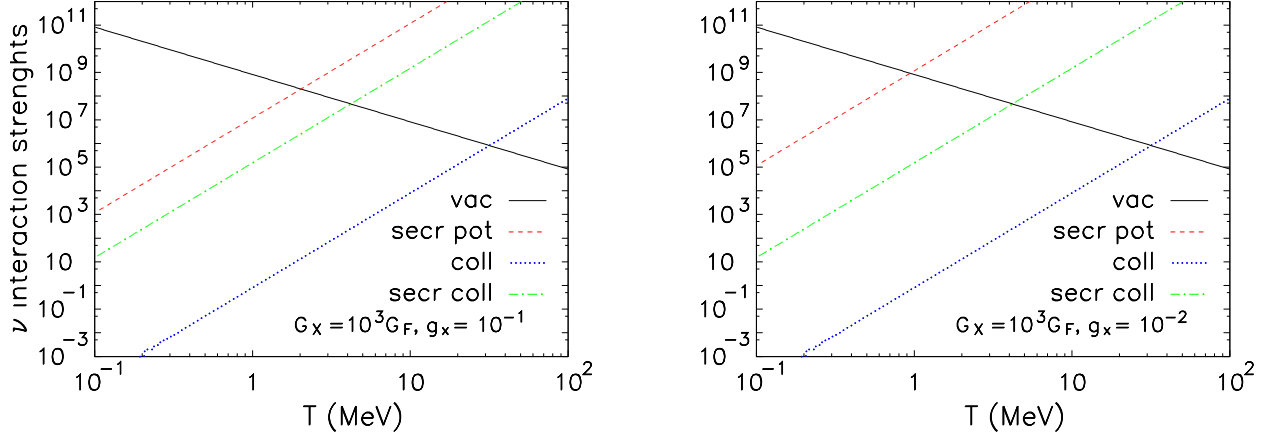


FIG. 1: Neutrino refractive and collisional rates (normalized in terms of the Hubble rate) versus temperature T for $G_X = 10^3 G_F$. Left panel corresponds to $g_X = 10^{-1}$, while right panel to $g_X = 10^{-2}$. The curves correspond to the active-sterile vacuum term (solid curve), the secret matter potential for $\rho_{ss} = 10^{-2}$ (dashed curve), the standard collisional term (dotted curve) and the collisional term associated with G_X^2 (dot-dashed curve).

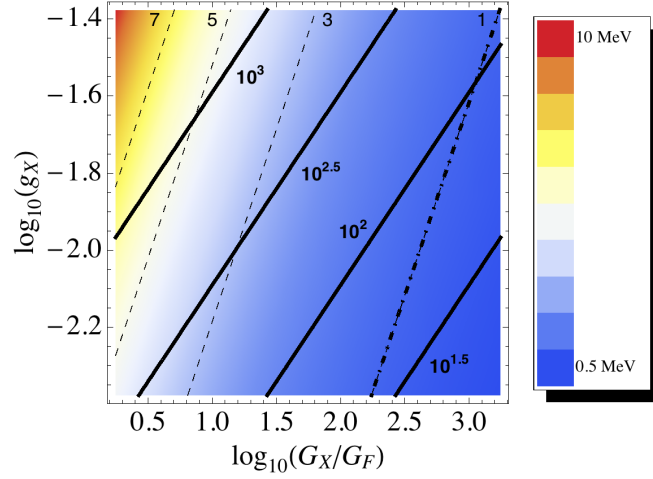


FIG. 2: Resonance temperature T_{res} in the plane (G_X, g_X) . Dashed curves represent constant T_{res} contours, while on solid curves M_X is constant. The values shown for both parameters are expressed in MeV. The case of $T_{\text{res}} = 1$ MeV is highlighted with a tick dot-dashed style, and correspond to the order of magnitude of BBN onset in the standard case, when neutron to proton density ratio freezes out.

Results are shown for $G_X = 10^3 G_F$. In the left panel we consider $g_X = 10^{-1}$ while in the right panel we take $g_X = 10^{-2}$, corresponding to $M_X = 390$ MeV and $M_X = 39$ MeV, respectively. In these plots we show the active-sterile vacuum term (solid curve) and the secret matter potential assuming $\rho_{ss} = 10^{-2}$ (dashed curve),

an indicative value when steriles are about to be excited. We realize that for $g_X = 10^{-1}$ the active-sterile resonance occurs at $T \simeq 2$ MeV, while for $g_X = 10^{-2}$ it takes place at $T \simeq 1$ MeV. We also show the standard collisional term (dotted) and the collisional term associated with G_X^2 (dot-dashed curve). Notice that in the active sector

the system remains collisional down to a few MeV, when the standard collision rate over Hubble rate drops below unity. On the other hand, the secret collisional term associated with G_X^2 remains larger than the Hubble rate till $T \sim 0.1$ MeV, implying that it will tend to drive the active-sterile system towards the flavor equilibrium for all the relevant evolution. The secret collisional term also dominates over the vacuum oscillation one for $T \gtrsim$ few MeV, thus breaking the coherence between different neutrino flavors and preventing significant oscillations at high temperatures.

In Fig. 2 we show the resonance temperature T_{res} for $\rho_{ss} = 10^{-2}$ in the plane (G_X, g_X) . From the resonance condition, one gets

$$\frac{T_{\text{res}}}{\text{MeV}} \simeq 0.4 \left[\left(\frac{\Delta m_{\text{st}}^2}{\text{eV}^2} \right) \left(\frac{M_X^2}{\text{MeV}^2} \right) \left(\frac{G_F}{G_X} \right) \frac{1}{\rho_{ss}} \right]^{1/6}. \quad (12)$$

Dashed and solid curves represent locations of constant T_{res} and M_X , respectively. We see that for the values of the parameters of the secret interactions considered in this figure, resonances may take place in a range of temperatures relevant for BBN. In the following we will focus on this range of parameters for our analysis.

We now consider the temperature evolution of neutrino momentum-averaged density matrix $\rho(T)$. In Fig. 3 we show the flavor evolution for the two cases of Fig. 1 compared to the standard case without secret interactions. In particular, in the left panels we represent ρ_{ee} (continuous curve), $\rho_{\mu\mu}$ (dotted curve) and ρ_{ss} (dashed curve) while right panels show the corresponding N_{eff} . The upper panel is the standard evolution in absence of sterile self-interactions; in this case sterile neutrinos are produced efficiently at $T \simeq 70$ MeV and they thermalize, giving $N_{\text{eff}} = 4$. Since the sterile production occurs at temperatures for which the active neutrinos are in a collisional regime, the latter are efficiently repopulated with $\rho_{ee} = \rho_{\mu\mu} = 1$. In the presence of secret ν_s - ν_s interactions, the sterile neutrino production is suppressed with respect to the standard case, till a resonance is intercepted. For the case of $g_X = 10^{-1}$ the production starts around $T = 3$ MeV. Due to a strong collisional term in the sterile neutrino sector, the system tends to evolve towards a flavor equilibrium with equal density of actives and sterile species. In this case since the flavor conversions occur in a temperature range where the active neutrinos are still in a collisional regime, though close to decoupling, this effect tends to repopulate the active sector, producing a final $\rho_{ee} \simeq 0.9$ and $\Delta N_{\text{eff}} = 0.8$. Finally, for $g_X = 10^{-2}$ active-sterile neutrino conversions occur at $T \lesssim 1$ MeV, where active neutrinos are no longer repopulated by the collisional effects. In this case, the depletion of the ρ_{ee} is remarkable, with a final value $\rho_{ee} = 0.7$ and $\Delta N_{\text{eff}} = 0.18$.

We have performed a scan of neutrino evolution as function of G_X and g_X focusing, as already mentioned, on the range where sterile productions is expected may alter the results of BBN. Our findings are summarized in Fig. 4, where we report the asymptotic values of ΔN_{eff}

in the (G_X, g_X) plane. Some reference values are shown as dashed lines.

III. IMPACT ON OBSERVABLES: THE LIGHT NUCLEI ABUNDANCES

A. Data and analysis

As known, BBN proceed in two basic steps. At the MeV scale weak processes maintaining chemical equilibrium between neutrons and protons become ineffective, when their rates drop below the value of the Hubble expansion rate. Later on, at $T \sim 80$ keV deuterium forms, and soon the whole nuclear reaction chain starts till it eventually stops at a temperature of order 10 keV. The role of neutrinos, including possible sterile states, is twofold. They gravitate and contribute as relativistic species to the total energy budget. This effect is encoded into a single parameter, N_{eff} . In addition, electron neutrino distribution in phase space is a key input for the weak proton-neutron rates, which fix the n/p ratio.

In the following we will exploit both the abundances of ^4He , typically cast in terms of its mass fraction Y_p , and the deuterium to hydrogen number density ratio, $^2\text{H}/\text{H}$. In fact, for both these nuclei we have trustable experimental determinations of their primordial values, which we will very briefly discuss later.

We start by summarizing how typically non standard neutrino abundances influence final light nuclei yields. For a fixed baryon density an increase of ΔN_{eff} shifts the weak rate freeze out at larger temperatures, since the Hubble rate is proportionally larger. From chemical equilibrium, which we can trust down to the freezing temperature T_{fr} , $n/p \sim \exp(-\Delta m/T_{\text{fr}})$, with Δm the neutron-proton mass difference. If T_{fr} gets higher due to an increase of N_{eff} , there are more neutrons available at the onset of deuterium formation, and this translates into a larger value of Y_p . This is what we expect if sterile neutrino states are excited in the early universe. This effect can be compensated if electron neutrino number density is *larger*, since in this case weak rates get increased. The secret interaction scenario, while may reduce the value of ΔN_{eff} to some (positive) values smaller than unity yet, it typically leads to a *smaller* electron neutrino density $\rho_{ee} < 1$. These two features, i.e. a positive ΔN_{eff} and less electron neutrinos, both conspire to produce a larger Y_p .

The way deuterium changes is less obvious. For a larger N_{eff} but keeping constant the ν_e distribution, the ratio $^2\text{H}/\text{H}$ increases. Nevertheless, in the case at hand it is quite difficult to get a simple feeling of how it would change depending on the two parameters G_X and g_X .

To exploit in a quantitative manner the BBN predictions for light nuclei, we have implemented the secret interaction scenario in the public numerical code **ParthENoPE** [46]. The value of ΔN_{eff} obtained by solving the equations of motion in Eq. (4) as function of the photon temperature is now read as an external input

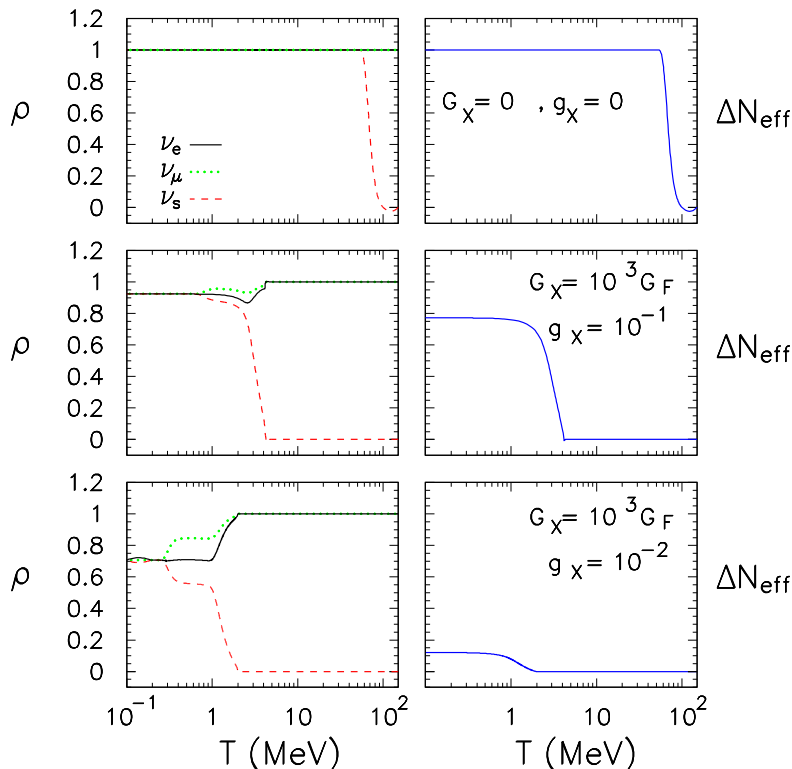


FIG. 3: Flavor evolution as functions of temperature T for different cases for $G_X = 10^3 G_F$. Upper panel is the standard case without secret interactions. Middle and lower plots are for $g_X = 10^{-1}$ and $g_X = 10^{-2}$, respectively. In the left panels we report ρ_{ee} (continuous curve), $\rho_{\mu\mu}$ (dotted curve) and ρ_{ss} (dashed curve). Right panels show the corresponding ΔN_{eff} .

by **ParthENoPE**, as well as the electron neutrino distribution which is used to compute the weak thermal rates of the neutron/proton reactions. We recall that neutrino distributions are evolved in the average-momentum approximation, so the ν_e distribution used in the thermal rate is the standard Fermi-Dirac function times $\rho_{ee}(T)$. Actually, at the level of approximation we are interested in, we compute the tree level Born weak rates with the modified electron neutrino distribution instead of using the standard rates employed in the public version of the code, which also contains the contribution of radiative corrections. The latter are quite involved function depending also on neutrino distribution, which should be in principle recomputed. The approximation we are using is thus to assume that the effect of a different ν_e distribution in the one-loop contribution rescales in the same way the Born rates do. We recall that radiative corrections

contribute typically for 4-5 % of the Born rates around the freeze out temperature, see e.g. [47]. Therefore, we neglect them.

In the analysis we have used the last updated result on Y_p reported in [48], based on a regression to zero metallicity with new He I emissivities, and using the dataset of [49]

$$Y_p = 0.2465 \pm 0.0097, \quad (13)$$

while for deuterium we consider the recent result of [21]

$$^2\text{H}/\text{H} = (2.53 \pm 0.04) \times 10^{-5}. \quad (14)$$

This value is the result of a reanalysis of all known deuterium absorption systems, including the new discovered very metal-poor damped Lyman- α system at redshift $z = 3.06726$ toward the QSO SDSS J1358+6522. Notice the quite small uncertainty, of the order of 1.6 %.

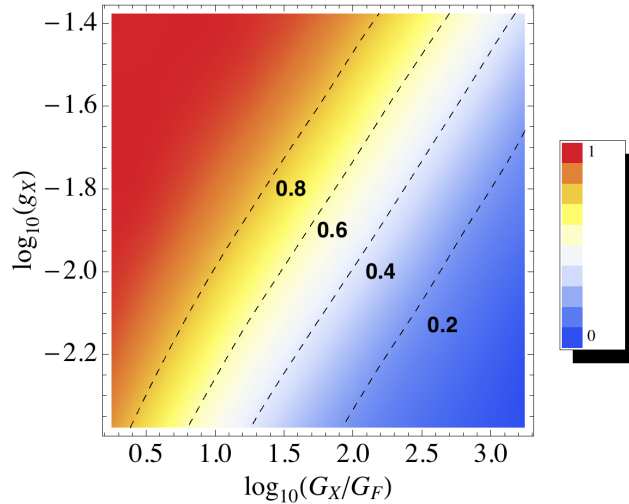


FIG. 4: The asymptotic values of ΔN_{eff} versus G_X and g_X . Colours from blue (lower right corner) to red (upper left corner) correspond to increasing values. Dashed curves show some reference values.

Before discussing our findings we make a last remark. While the uncertainty on the theoretical value of ^4He from **PARthENoPE** is extremely small⁴ and negligible with respect to the experimental uncertainty of Eq. (13), the prediction for deuterium abundance is still affected by a large error (mainly) due to the present uncertainty on the astrophysical factor of the $d(p, \gamma)^3\text{He}$ reaction, which is the leading destruction channel of this nuclide. Once we propagate the uncertainty reported in [51], see also [52], the value of $^2\text{H}/\text{H}$ change by the amount $\pm 0.06 \times 10^{-5}$, which is even larger than the experimental error of Eq. (14). A theoretical *ab initio* calculation of the S-factor for this process is also available, which suggests a higher cross section in the center of mass energy range relevant for BBN [53–55], whose impact has been recently analyzed in [56, 57]. The predicted value of deuterium is in this case lower than if we used the experimental best fit of the rate, and in a better agreement with the experimental result of [21]. In order to have a clear assessment of the error budget on deuterium theoretical prediction, we have decided to conservatively use the present experimental results of [51], but we stress that it would be important also for the issue considered in this paper to have new experimental data on $d(p, \gamma)^3\text{He}$ in the BBN energy range, with a higher precision. If the theoretical *ab initio* calculation would be confirmed, and the experimental error on the astrophysical factor of the $d(p, \gamma)^3\text{He}$ reaction would be reduced by say, a factor three, which seems plausible [57], the constraints we will discuss later would become more stringent.

B. Results

Our results are summarized in Figures 5 and 6, showing the bounds on the secret interaction parameters (G_X, g_X) coming from Y_p and deuterium, respectively. By looking at these panels one can grasp the typical dependence of $^2\text{H}/\text{H}$ and Y_p upon the two relevant parameters. Namely, helium is an *increasing* function of G_X for fixed g_X , while it decreases with g_X for a given G_X . In the same range deuterium shows exactly the opposite behavior. Consider first the helium mass fraction Y_p . The two dark regions in Fig. 5 in the upper left and lower right corners are the allowed G_X and g_X at 1.5σ , where σ is the experimental error on Y_p of Eq. (13), the theoretical error being completely negligible. The value of the baryon density is varied in the 95 % C.L. region of Planck results, $\Omega_B h^2 = 0.02207 \pm 0.00054$ [58]. The solid line bounds instead the 1.5σ allowed region if we fix $\Omega_B h^2$ at the Planck best fit value. We see that for both these cases almost all the parameter space is excluded. We have shown the 1.5σ exclusion contours, since for the 2σ cases the whole range for G_X and g_X of Fig. 5 *would have been permitted*. Conversely, at 1σ the whole plane would be instead *excluded*. In fact, the uncertainty on Y_p , at the level of 4 %, is too large to severely constrain the secret interaction parameter, which are indeed all permitted at 2σ . Instead, as we will show later, the deuterium constraints will be much stronger.

In order to understand the dependence of Y_p on the secret interaction parameters, note that if we decrease g_X for a given G_X , the electron neutrino distribution becomes smaller, when the production takes place at smaller temperatures, where electron neutrinos are less efficiently repopulated by pair creation processes (see e.g. Fig. 2). It follows that also weak rates decrease. At the same time, decreasing g_X also ΔN_{eff} becomes smaller and

⁴ The only source of uncertainty is in fact, the small error on neutron lifetime, $\tau_n = 880.0 \pm 0.9$ sec [50].

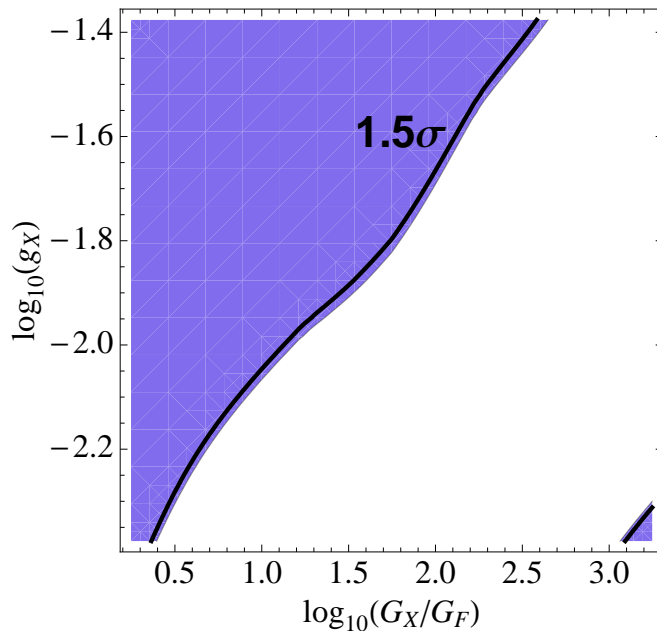


FIG. 5: ^4He results. The dark region is the 1.5σ allowed parameter space for the helium mass fraction Y_p using the experimental result of Eq. (13), varying the baryon density parameter in the range $0.02153 \leq \Omega_B h^2 \leq 0.02261$, corresponding to the 95 % C.L. Planck range. The solid line bounds the permitted region if we fix $\Omega_B h^2 = 0.02207$, the best fit quoted by Planck collaboration. At 2σ the whole region shown for G_X and g_X would be allowed, while at 1σ it is all excluded.

thus the expansion rate is also lowered. The combination of two effects causes an increase in Y_p when lowering g_X . To get the same value of Y_p it is thus, necessary to decrease the value of G_X , and this explain the bending of the iso-helium contours, as the one depicted in the upper left corner of Fig. 5.

These considerations only apply as long as the resonance temperature is larger or of the order of the typical scale of n/p ratio freezing. If sterile neutrinos are produced below this scale and electron neutrino distribution is still almost unchanged at say, $T \simeq 0.8$ MeV, BBN start being insensitive to the presence of sterile secret interactions. This explains why we can see another branch of allowed values for G_X and g_X in the lower right corner of Fig. 5: in this region resonance temperature is of order 0.6–0.7 MeV, and Y_p is less affected.

We discuss now the bound from $^2\text{H}/\text{H}$. In Fig. 6 the regions below the various curves are the allowed ranges for the two parameters at a number of σ reported for each curve. With σ here we denote the total error obtained by summing in quadrature the experimental error of Eq. (14) and the one due to nuclear rate uncertainties [51]. From the left panel of the Figure we see that if baryon density is fixed at the Planck best fit, almost all the range shown for G_X and g_X is excluded at 3σ , but the high- G_X and small- g_X lower right corner. If we translate the bound in terms of the X boson mass, we obtain at 3σ as allowed mass range

$$M_X \leq 40 \text{ MeV} . \quad (15)$$

Notice that in absence of sterile neutrinos the

ParthEnoPE code prediction for $^2\text{H}/\text{H}$ at $\Omega_B h^2 = 0.02207$ is $^2\text{H}/\text{H} = (2.65 \pm 0.07) \times 10^{-5}$, which is already larger than the experimental value we are using, and compatible with it at the 2σ level. When switching on the ν_s secret interactions, deuterium relic abundance is strongly enhanced, unless we consider the low- g_X and high- G_X range, the lower-right corner of Fig. 6, where now the agreement with experimental value can be at least at the level of 3σ . The improvement in this region of the agreement between the theoretical values of deuterium and the experimental determination of Eq. (14) is due to the fact that in this range of parameters the resonance temperature is below or of the order of 0.8 MeV (see Fig. 2), the typical scale of n/p ratio freezing. As already noticed, in this limit BBN starts being blind to electron neutrino distortion and the only effect is the faster expansion rate due to a positive ΔN_{eff} , which gives a slightly more deuterium abundance.

The constraint from $^2\text{H}/\text{H}$ is weaker if we take a larger value of $\Omega_B h^2 = 0.02261$, the 95 % C.L. upper limit from Planck and the bound on the X boson mass becomes $M_X \leq 220$ MeV at 3σ . This is because deuterium is a rapidly decreasing function of baryon fraction, so one way to compensate for a higher theoretical prediction is to shift $\Omega_B h^2$ towards larger values. In this case too, however, most of the parameter range is excluded at 2σ . We conclude that the secret interaction scenario is in tension with present data on primordial deuterium, unless baryon density is quite at the upper boundary of Planck result.

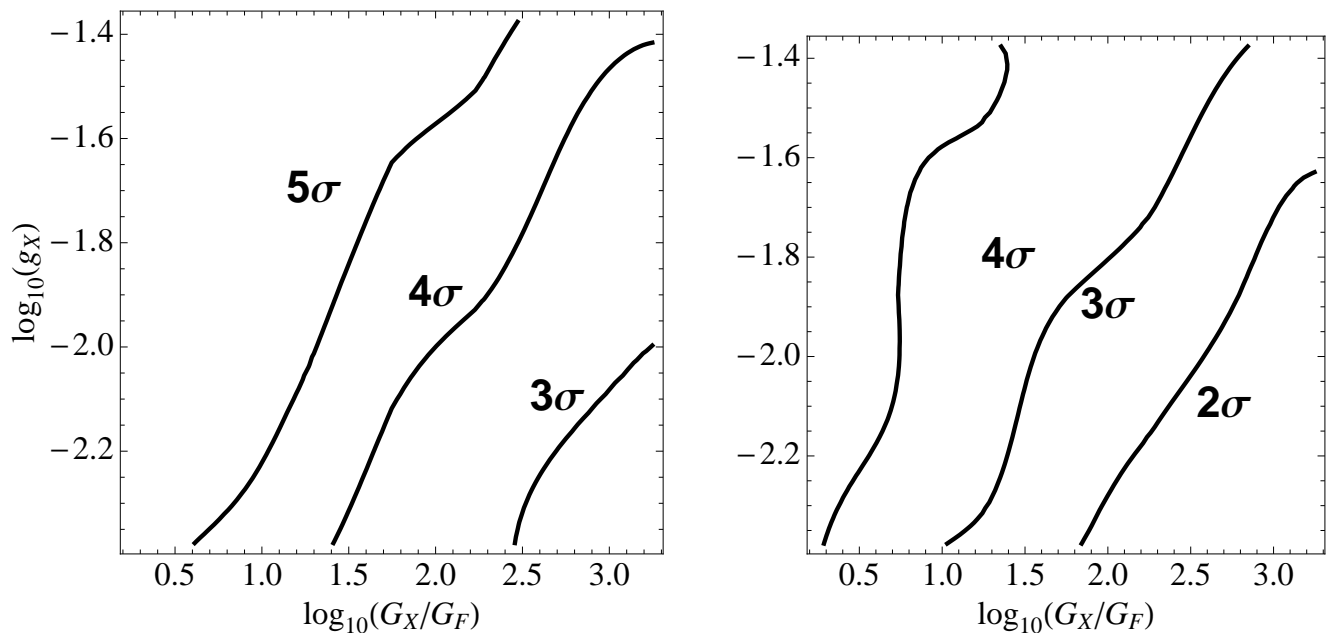


FIG. 6: $^2\text{H}/\text{H}$ results. Left panel: the region below each curve is the allowed one at the number of σ shown for each case, using the experimental determination for $^2\text{H}/\text{H}$ of Eq. (14) and for $\Omega_B h^2 = 0.02207$, the best fit quoted by Planck collaboration. Right panel: the same as in the left panel but for a higher baryon density $\Omega_B h^2 = 0.02261$, i.e. the Planck upper limit at 95 % C.L.

IV. CONCLUSIONS

Secret interactions among sterile neutrinos, mediated by a gauge boson with $M_X \ll M_W$, have been recently proposed [26, 27] as a possible mechanism to suppress the thermalization of eV sterile neutrinos in the early universe, by the large matter potential they generate in this case. However, the active-sterile neutrino ensemble would also experience a resonance when the matter term gets close to the vacuum oscillation frequency. In this situation the sterile neutrino production would be enhanced by the combination of a resonant production, with a non-resonant one due to the large collisional effects, caused by the secret interactions in the sterile ν sector.

For values of the coupling constant $g_X \gtrsim 10^{-2}$ and masses of the gauge boson $M_X \gtrsim \mathcal{O}(10)$ MeV [26], the sterile neutrino production would occur at epochs relevant for the BBN. In this paper we have analyzed the BBN bounds on the secret interaction scenario. The standard BBN dynamics is indeed, possibly changed by a larger value of the number of effective relativistic degrees of freedom, N_{eff} , and the spectral distortions on ν_e induced by the active-sterile flavor conversions. Using the present determination of ^4He and deuterium primordial abundances, we found that due to the 4 % error on experimental determination of helium mass fraction, Y_p gives no significant bounds. We comment that very recently a new measurement of $Y_p = 0.2551 \pm 0.0022$ has been presented [59]. The smaller value of the quoted error would allow one to put stronger limits than the ones we obtained.

The $^2\text{H}/\text{H}$ density ratio excludes much of the parameter space if one assume a baryon density at the best fit value of Planck experiment, $\Omega_B h^2 = 0.02207$. In this case we can set an upper limit on the X boson mass at 3σ $M_X \leq 40$ MeV. This bound becomes weaker for a higher baryon fraction, $\Omega_B h^2 = 0.02261$, which is the 95 % C.L. upper bound of Planck, that is $M_X \leq 220$ MeV. The bound on M_X can be improved measuring the astrophysical factor of the $d(p, \gamma)^3\text{He}$ process cross section in the relevant energy range, with a smaller uncertainty. Therefore, new experiments measuring this quantity are therefore very welcome.

As a consequence of our analysis the parameter space for secret interactions to reconcile sterile neutrinos with cosmology is significantly reduced. A possible way to avoid the BBN bounds is to choose the mass of the mediator so light as $M_X \lesssim 1$ MeV to shift the resonances at temperatures too low ($T \lesssim 0.1$ MeV) to be relevant for nucleosynthesis, as considered in [8]. However, in this situation the sterile neutrino production may have an effect on the cosmological neutrino mass bound [60]. Results of this study will be presented elsewhere [61].

Acknowledgements

We thank Silvia Pascoli for interesting discussions. N.S. also acknowledges Steen Hannestad and Thomas Tram for valuable clarifications and suggestions. G. M., and O.P. acknowledge support by the *Istituto Nazionale di Fisica Nucleare* I.S. FA51 and the PRIN 2010 “Fisica

Astroparticellare: Neutrini ed Universo Primordiale” of the Italian *Ministero dell’Istruzione, Università e Ricerca*. The work of A.M. was supported by the German Science Foundation (DFG) within the Collaborative

Research Center 676 “Particles, Strings and the Early Universe.” N.S. acknowledges support from the European Union FP7 ITN INVISIBLES (Marie Curie Actions, PITN- GA-2011- 289442).

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